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# Acoustic Instabilities as a Source of Turbulence in Reduced Density Channels

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*Plasma Physics Division*

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# ACOUSTIC INSTABILITIES AS A SOURCE OF TURBULENCE IN REDUCED DENSITY CHANNELS

## I. Introduction

In experiments at NRL hot, reduced density channels have been formed in air and argon by ohmic heating<sup>1</sup> and in nitrogen by laser heating.<sup>2</sup> Nominally such channels could persist, in pressure equilibrium with the cold surrounding gas, until they cooled by heat conduction. In fact they become turbulent and dissipate  $\sim 2000$  times faster than heat conduction would dictate. Boris and Picone<sup>3</sup> have shown that spatial asymmetries in the initial heating of the channels can cause such turbulence. However the appearance of turbulence in channels with relatively symmetric initial heating suggests the possibility of additional sources.

Electrical measurements on the channels in air show that they do not decay to an equilibrium electron density for  $\sim 100 \mu s$ . In argon significant amounts of energy can be trapped (resonant self trapping) and concentrated in metastable levels. Air and nitrogen can also store energy in vibrational states with long relaxation times. It has been recognized that acoustic wave growth can occur in a gas maintained in a nonequilibrium state.<sup>4</sup> This paper demonstrates that such growth can occur in a gas relaxing from a nonequilibrium condition. The theory is then applied to the channel situation to see if energy "frozen in" during the channel expansion is sufficient to cause a disturbance to amplify to the point of becoming turbulent.

Finally we show the channels are stable, if the gas remains in thermal equilibrium, despite the gas possessing a temperature dependent specific heat.

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## II. A Description of the Acoustic Instability

As a simple model which will demonstrate the growth phenomena we consider a diatomic gas possessing a long lived vibrational level. We suppose this level may be depopulated only by binary collisions whereupon its energy is given up to the internal kinetic energy. If the gas properties remain spatially uniform a slow relaxation to thermal equilibrium occurs. A sound wave passing through the gas causes local perturbations in the temperature and density and thus local perturbations in the relaxation rate. A region of compression speeds the relaxation rate delivering an extra energy "kick" to the subsequent re-expansion. On the other hand slowing of the relaxation rate in the rarefactions assists their re-collapse. The sound wave thus grows by diverting some of the energy from the decaying vibrational state into bulk motion and away from internal kinetic energy.

Sound waves are normally damped by viscosity and by heat conduction between the hot, compressed and cold, rarefied zones. We will neglect viscosity and heat conduction in the present analysis. This approximation must thus be justified, in any application, by comparison of the predicted growth rate to the so-called classical absorption.<sup>5</sup>

Finite relaxation times for internal degrees of freedom normally introduce additional damping due to the loss and retrieval of energy to and from these states. This effect is outweighed in the present case where the vibrational level is too high lying to be accessible at the post expansion channel temperature.

## III. Linear Stability Analysis

We describe the behavior of the gas by the usual Eulerian variables density,  $\rho$ , pressure,  $p$ , and velocity,  $v$ , plus a specific internal energy exclusive of the vibrational level,  $\epsilon$ , and a specific energy in the vibrational level alone,  $\epsilon^*$ . The latter energy is related to the density of excited atoms,  $n^*$ , via

$$\epsilon^* = \frac{En^*}{\rho} \quad (1)$$

where  $E$  is the excitation energy. The rate equation determining  $n^*$

$$\frac{\partial n^*}{\partial t} = -\sigma n^* n \bar{v} + \sigma^* n (n - n^*) \bar{v} \quad (2)$$

may be expressed in terms of  $\epsilon^*$ ,  $\rho$ , and  $\epsilon$ . (We neglect the last term representing relaxation damping since we assume  $\sigma^* \ll \sigma$  and  $n^* \geq n$ .)

$$\frac{\partial \epsilon^*}{\partial t} = -\sigma \epsilon^* \rho \sqrt{\frac{32 A_0 \epsilon}{5 \pi A m}} \quad (3)$$

where

$\sigma^* \approx$  cross section for collisional excitation

$\sigma \approx$  cross section for collisional de-excitation

$n \approx$  total number density

$\bar{v} \approx$  mean velocity between identical

$$\text{molecules} = \sqrt{2} \sqrt{\frac{8kT}{\pi m}}$$

$A_0 \approx$  Avogadro's number

$A \approx$  gram molecular weight

$m \approx$  mass per molecule

Equation (3) may be rewritten for a moving fluid by following standard procedures<sup>6</sup>

$$\frac{D\epsilon^*}{Dt} = A \epsilon^* \rho \sqrt{\epsilon} \quad (\text{rate equation}) \quad (4)$$

where  $D/Dt$  is the material derivative given by

$$\frac{D}{Dt} = \frac{\partial}{\partial t} + \mathbf{v} \cdot \nabla$$

and we have defined

$$A \equiv \sigma \sqrt{\frac{32 A_0}{5 \pi A m}}$$

Other governing equations in addition to (4) include

$$\frac{D\rho}{Dt} = -\rho \nabla \cdot \mathbf{v} \quad (\text{conservation of mass}) \quad (5)$$

and

$$\rho \frac{D\mathbf{v}}{Dt} = -\nabla p \quad (\text{conservation of momentum}) \quad (6)$$

where viscosity has been neglected. The conservation of energy equation contains an extra term to reflect the transfer of energy from the vibrational level into internal kinetic energy

$$\rho \frac{D\epsilon}{Dt} = -\rho \nabla \cdot \mathbf{v} - \rho \frac{D\epsilon^*}{Dt} \quad (\text{conservation of energy}) \quad (7)$$

and heat conduction has been neglected.

One additional equation is necessary to provide a closed set. Since the vibrational energy interacts with the fluid only via the decay represented by the last term in (7), the fluid variables are interrelated as in an equilibrium gas. We may therefore adopt the usual equation of state<sup>7</sup>

$$\epsilon = \frac{1}{\gamma - 1} \left( \frac{p}{\rho} \right) \quad (\text{equation of state}). \quad (8)$$

A slowly relaxing, spatially uniform medium represents a succession of equilibrium states to a sound wave whose period is short compared to the decay time. The fluid variables may be represented as a slowly changing quasi-equilibrium value plus a rapidly changing perturbation

$$\epsilon^* = \epsilon_0^* + \epsilon_1^* \quad \text{etc.} \quad (9)$$

Two possible types of quasi-equilibria are of interest, constant pressure,  $p_0$ , and constant density,  $\rho_0$ . The former more nearly represents the situation within the channel but either is acceptable under the slow relaxation approximation. The latter leads to simpler algebra and we shall develop that case. Our assumption of a constant density, spatially uniform equilibrium implies

$$\frac{\partial \rho_0}{\partial t} = 0 \quad (10)$$

$$\nabla \rho_0 = \nabla \epsilon_0 = \nabla \epsilon_0^* = 0. \quad (11)$$

These assumptions, applied to the governing equations, (4) through (8), show that the equilibrium is further described by

$$\frac{\partial \epsilon_0^*}{\partial t} = -4 \epsilon_0^* \rho_0 \sqrt{\epsilon_0} \quad (12)$$

$$\mathbf{v}_0 = 0 \quad (13)$$

$$\nabla p_0 = 0 \quad (14)$$

$$\frac{\partial \epsilon_0}{\partial t} = - \frac{\partial \epsilon_0^*}{\partial t} \quad (15)$$

$$\epsilon_0 = \frac{1}{\gamma - 1} \left( \frac{p_0}{\rho_0} \right). \quad (16)$$

If we substitute the split forms of the variables, (9), into the governing equations, (4) through (8), eliminate those terms which describe the equilibrium behavior by invoking (10) through (16), and retain only terms to first order in the perturbations, we are lead to the linearized equations for the perturbation variables

$$\frac{\partial \epsilon_1^*}{\partial t} = -A \left[ \rho_0 \epsilon_1^* \sqrt{\epsilon_0} + \epsilon_0 \dot{\rho}_1 \sqrt{\epsilon_0} + \frac{\epsilon_0 \dot{\rho}_0 \epsilon_1}{2\sqrt{\epsilon_0}} \right] \quad (17)$$

$$\frac{\partial \rho_1}{\partial t} = -\rho_0 \nabla \cdot \mathbf{v}_1 \quad (18)$$

$$\rho_0 \frac{\partial \mathbf{v}_1}{\partial t} = -\nabla p_1 \quad (19)$$

$$\rho_0 \frac{\partial \epsilon_1}{\partial t} = -\rho_0 \nabla \cdot \mathbf{v}_1 - \rho_0 \frac{\partial \epsilon_1^*}{\partial t} \quad (20)$$

$$\epsilon_1 = \epsilon_0 \left( \frac{p_1}{p_0} - \frac{\rho_1}{\rho_0} \right). \quad (21)$$

We need only consider planar disturbances to demonstrate the growth phenomena. Given linear equations it is sufficient to consider separately each Fourier component of any sonic disturbance. Substituting a normal mode form for each perturbation variable

$$\epsilon_1^* = \epsilon_1^*(\omega, k) e^{i(kx - \omega t)} \quad \text{etc.} \quad (22)$$

into (17) through (21) yields

$$i\omega \epsilon_1^* = -A \left[ \rho_0 \epsilon_1^* \sqrt{\epsilon_0} + \epsilon_0 \dot{\rho}_1 \sqrt{\epsilon_0} + \frac{\epsilon_0 \dot{\rho}_0 \epsilon_1}{2\sqrt{\epsilon_0}} \right] \quad (23)$$

$$\omega \rho_1 = k \rho_0 v_1 \quad (24)$$

$$\omega \rho_0 v_1 = k p_1 \quad (25)$$

$$\omega \rho_0 \epsilon_1 = k \rho_0 v_1 - \omega \rho_0 \epsilon_1^* \quad (26)$$

$$\epsilon_1 = \epsilon_0 \left( \frac{p_1}{p_0} - \frac{\rho_1}{\rho_0} \right). \quad (27)$$

The perturbation variables are now understood to be functions of  $\omega$  and  $k$ .

Equations (23) through (27) may be solved for any perturbation variable, the pressure proves to be a convenient choice. We solve (23) for  $\epsilon_1^*$ , and use that result, and (27), to eliminate  $\epsilon_1$  and  $\epsilon_1^*$  in (26). Equations (24) and (25) may be used to eliminate  $\rho_1$  and  $v_1$  resulting in

$$\rho_1 \left( \frac{\omega^2 \rho_0 \epsilon_0}{\rho_0} + \frac{\omega^2 \rho_0^2 A \epsilon_0^* \sqrt{\epsilon_0}}{2 \rho_0 (i\omega - A \rho_0 \sqrt{\epsilon_0})} - \epsilon_0 k^2 - \frac{\rho_0}{\rho_0} k^2 + \frac{\rho_0 A \epsilon_0^* k^2 \sqrt{\epsilon_0}}{2 (i\omega - A \rho_0 \sqrt{\epsilon_0})} \right) = 0. \quad (28)$$

Setting the bracketed term to zero yields the dispersion relation between  $\omega$  and  $k$ . This becomes, upon further simplification

$$\omega^2 (2i\omega \sqrt{\epsilon_0} - 2\epsilon_0 A \rho_0 + \rho_0 A \epsilon_0^*) - k^2 (2ic^2 \omega \sqrt{\epsilon_0} - 2c^2 A \rho_0 \epsilon_0 - A \rho_0 \epsilon_0^*) = 0 \quad (29)$$

where we have introduced the equilibrium sound speed

$$c^2 = \frac{\rho_0}{\rho_0 \epsilon_0} \left( \epsilon_0 + \frac{\rho_0}{\rho_0} \right) = \frac{\gamma \rho_0}{\rho_0} \quad (30)$$

(Notice that turning off the coupling between  $\epsilon$  and  $\epsilon^*$ , by setting  $A$  to zero recovers the dispersion relation for normal sound waves,  $\omega^2 = c^2 k^2$ .)

The solution of equation (29) requires  $\omega$  to be complex. We may substitute  $\omega = \omega_r + i\omega_i$  in (29) and separately equate the real and imaginary parts of the resulting expression to zero. This gives

$$\begin{aligned} & \omega_r^2 (6\omega_r \sqrt{\epsilon_0} + 2\epsilon_0 A \rho_0 - \rho_0 A \epsilon_0^*) \\ & - k^2 (2c^2 \omega_r \sqrt{\epsilon_0} + 2c^2 A \rho_0 \epsilon_0 + A \rho_0 \epsilon_0^*) \\ & - \omega_i^2 (2\omega_r \sqrt{\epsilon_0} + 2\epsilon_0 A \rho_0 - \rho_0 A \epsilon_0^*) = 0 \end{aligned} \quad (31)$$

and

$$3\omega_r^2 \sqrt{\epsilon_0} + (2\epsilon_0 A \rho_0 - \rho_0 A \epsilon_0^*) \omega_r + \sqrt{\epsilon_0} (c^2 k^2 - \omega_r^2) = 0. \quad (32)$$

Equation (31) may be used to eliminate  $\omega_i^2$  in (32) with the result that

$$\omega_r^3 + M\omega_r^2 + N\omega_r + Q = 0 \quad (33)$$

where

$$M = A \rho_0 \sqrt{\epsilon_0} - \frac{A \rho_0 \epsilon_0^*}{2 \sqrt{\epsilon_0}} \quad (34)$$

$$N = \frac{\epsilon_0 \rho_0^2 A^2}{4} - \frac{\rho_0^2 A^2 \epsilon_0^*}{4} + \frac{\rho_0^2 A^2 \epsilon_0^{*2}}{16 \epsilon_0} + \frac{c^2 k^2}{4} \quad (35)$$

$$Q = - \frac{A k^2 \rho_0 \epsilon_0^* (\gamma - 1)}{16 \sqrt{\epsilon_0}} \quad (36)$$

The constant term,  $Q$ , will always be negative, thus there will always be a positive value of  $\omega$ , corresponding to wave growth.

#### IV. Application to Laser Formed Channels in Nitrogen

The laser formed channel in nitrogen reaches the following conditions after expanding to pressure equilibrium

$$\begin{aligned} p_0 &= p_{\text{ambient}} = 10^6 \frac{\text{dynes}}{\text{cm}^2} \\ \rho_0 &= \frac{\rho_{\text{ambient}}}{4.33} = 2.3 \times 10^{-4} \frac{\text{g}}{\text{cm}^3} \\ T &= \frac{\rho_{\text{ambient}} T_{\text{ambient}}}{\rho_0} = 1300 \text{ K} \\ \epsilon_0 &= \epsilon_{\text{trans.}} + \epsilon_{\text{rot.}} = \frac{5}{2} \frac{A_0}{A} kT = 10^{10} \frac{\text{erg}}{\text{g}}. \end{aligned}$$

The expansion is approximately adiabatic. If the vibrational energy is frozen in during this expansion the effective adiabatic exponent is  $\gamma \approx 1.4$  reflecting only the translational and rotational contributions to the specific heat. We may backtrack to find the initial temperature prior to expansion

$$T_i = T \left( \frac{\rho_{\text{ambient}}}{\rho_0} \right)^{\gamma-1} = 2400 \text{ K}. \quad (37)$$

The frozen approximation is justified by calculating the vibrational relaxation times ( $\tau_i$  and  $\tau$ ) under both the pre and post expansion conditions. The relaxation time<sup>8</sup> depends on the binary collision rate,  $\tau_{\text{coll.}}$ , and the probability of de-excitation as a result of a collision,  $p_{10}$ \*

$$\tau = \frac{\tau_{\text{coll.}}}{p_{10}}. \quad (38)$$

Noting that

$$\tau_{\text{coll.}} = \frac{1}{n \sigma_{\text{coll.}} \bar{v}} \quad (39)$$

we find

$$\begin{aligned} \tau_i &= 5 \mu \text{ sec} \\ \tau &= 300 \mu \text{ sec} . \end{aligned}$$

\*The de-excitation cross section,  $\sigma$ , introduced in Section III and  $p_{10}$  are related in via  $\sigma = p_{10} \sigma_{\text{coll.}}$ .

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The channel expansion takes  $\sim 30 \mu \text{ sec}$  so vibrational levels may first become excited and then frozen.

In nitrogen only the lowest vibrational level ( $E/k \approx 3340K$ ) will be appreciably excited under the initial conditions. As an upper limit on the possible vibrational energy we calculate the equilibrium population for this state

$$n_i^* \approx \frac{n_i e^{-E/kT_i}}{1 + e^{-E/kT_i}} \approx \frac{n_i}{5} \quad (40)$$

If this fraction of the molecules remains excited during the expansion we find

$$\epsilon_0^* = \left( \frac{n_i^*}{n_i} \right) kT_i \frac{A_0}{A} \approx 1.5 \times 10^9 \frac{\text{erg}}{g}$$

We now calculate the growth rate,  $\omega_i$ . Under these conditions only the linear term in (33) and the last term in (35) are significant. The growth rate becomes independent of wave number,  $k$ , and is given by

$$\omega_i \approx \frac{A \rho_0 (\gamma - 1) \epsilon_0^*}{4 \gamma \sqrt{\epsilon_0}} = 36 \text{ sec}^{-1} \quad (41)$$

where we have used (see Eq. 12)

$$A = \frac{1}{\tau \rho_0 \sqrt{\epsilon_0}} \quad (42)$$

to evaluate  $A$ .

This growth rate is clearly too small to have any effect over the  $300 \mu \text{ s}$  during which the nitrogen relaxes.

## V. Stability of Gases with Temperature Dependent Specific Heats

For low frequencies or fast relaxation rates the gas remains in thermal equilibrium as it adjusts to the disturbance. We may include the contributions from all the internal degrees of freedom in  $\epsilon$  and set  $\epsilon^*$  to zero. The specific energy now acquires a non linear temperature dependency

$$\epsilon \propto T^2 \quad (43)$$

where  $a \geq 1$ . If we combine (43) with the gas law,  $p = \rho RT$ , and linearize the resulting expression we find

$$\epsilon_1 = a \epsilon_0 \left( \frac{p_1}{p_0} - \frac{\rho_1}{\rho_0} \right). \quad (44)$$

The dispersion relation for perturbations from fixed equilibrium is then given by substituting  $\epsilon_0 = a \epsilon_1$  and  $\epsilon_0^* = 0$  in (29). Ordinary sound waves now result ( $\omega = ck$ ) with a modified sound speed

$$c = \sqrt{\frac{p_0}{a \rho_0} \left( a \epsilon_0 + \frac{p_0}{\rho_0} \right)} \quad (45)$$

The medium is now stable to any planar disturbance.

The energy principle<sup>9</sup> permits us to demonstrate stability in complex geometric situations without having to solve the equations of motion. We shall apply it to a column of hot gas, with a temperature dependent specific heat, embedded in a cold surrounding medium.

The Lagrangian density in a region of uniform pressure and density may be expressed as<sup>10</sup>

$$L = T - V = \frac{1}{2} \rho_0 \dot{\eta}^2 + \frac{1}{2} v_0 \left( \frac{\partial p}{\partial v} \right)_0 (\nabla \cdot \eta)^2 \quad (46)$$

where  $\eta$  is a small fluid displacement from equilibrium and  $v_0 = 1/\rho_0$  is the specific volume. Proceeding as in the derivation of (44) we find

$$\partial \epsilon = a \epsilon_0 \left( \frac{\partial p}{\partial \rho} + \frac{\partial v}{\partial \rho} \right) \quad (47)$$

But having neglected heat conduction, the first law of thermodynamics states

$$\partial \epsilon = - p_0 \partial v. \quad (48)$$

Equating (47) and (48) gives

$$\left( \frac{\partial p}{\partial v} \right)_0 = - \frac{p_0}{a \epsilon_0} \left( p_0 + \frac{a \epsilon_0}{v_0} \right) \quad (49)$$

so that

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$$L = \frac{1}{2} \rho_0 \dot{\eta}^2 - \frac{1}{2} \frac{\rho_0 v_0}{a \epsilon_0} \left( \rho_0 + \frac{a \epsilon_0}{v_0} \right) (\nabla \cdot \eta)^2. \quad (50)$$

Applying Lagrange's equation for continuous media<sup>11</sup> to (50) we recover the wave equation with the sound speed given by (45).

The energy principle states that the motion of the system will be bounded if the total potential energy for the system increases for any disturbance. The potential energy is

$$\begin{aligned} V = \iiint d^3x \quad V = & \frac{1}{2} \iiint_{\text{column}} d^3x \frac{\rho_0 v_h}{a \epsilon_h} \left( \rho_0 + \frac{a \epsilon_h}{v_h} \right) (\nabla \cdot \eta)^2 \\ & + \frac{1}{2} \iiint_{\text{outside}} d^3x \frac{\rho_0 v_c}{a \epsilon_c} \left( \rho_0 + \frac{a \epsilon_c}{v_c} \right) (\nabla \cdot \eta)^2 \end{aligned} \quad (51)$$

where the subscripts *h* and *c* refer to hot and cold respectively. The potential energy is positive for any compressive (i.e. sonic) disturbance and zero for any incompressible flow. These situations correspond to bounded oscillations and persistent but non amplifying flows.

## VI. Conclusions

An excess of nonequilibrium internal energy in a gas transforms the relaxational damping of sound waves into a growth phenomena. A model diatomic gas has been considered which possesses one vibrational level. The small signal growth rate for normal modes (plane waves) has been calculated from the rate equation and the equations of motion. The growth proves to be slight for the conditions that obtain within the laser heated reduced density channels. Thus it appears that air chemistry plays no role in creating the turbulence observed in these channels. Metastable energies an order of magnitude or more larger than the kinetic energy are required to produce a significant effect.

An energy principle analysis shows that the internal energy stored in rapidly relaxing states cannot cause instabilities i.e. a gas which remains in thermal equilibrium but possesses a temperature dependent specific heat does not constitute an amplifying medium for sound waves.

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